Chiral dynamics for exotic open-charm resonances

M.F.M. Lutz[1] and E.E. Kolomeitsev[2]

[1]*Gesellschaft für Schwerionenforschung (GSI)*
Planck Str. 1, 64291 Darmstadt, Germany
[2]*The Niels Bohr Institute, Blegdamsvej 17, DK-2100 Copenhagen, Denmark*

We review the open-charm spectrum of chiral excitations in QCD. At leading order in a chiral expansion a parameter-free prediction is obtained for the coupled-channel scattering of Goldstone bosons off open-charm mesons and baryons with $0^-, 1^-$ and $\frac{1}{2}^+$ quantum numbers. The recently announced narrow $D_s(2317)$ and $D_s(2463)$ mesons observed by the BABAR and CLEO collaborations are reproduced. Also the baryon states $\Lambda_c(2593), \Lambda_c(2880)$ and $\Xi_c(2790)$ discovered earlier by the CLEO collaboration are recovered naturally. Besides describing resonances with conventional quantum numbers additional multiplets with members that carry exotic quantum numbers are predicted. In particular, so far unobserved narrow isospin-singlet open-charm mesons with negative strangeness are generated dynamically. Similarly, narrow isospin-doublet open-charm baryons with positive strangeness are suggested.

I. INTRODUCTION

Recently a new narrow state of mass 2.317 GeV that decays strongly into $D_s^+ \pi^0$ was announced [1]. This result was confirmed [2] and a second narrow state of mass 2.463 GeV decaying into $D_s^* \pi^0$ was observed. Such states were first predicted in [3, 4] based on the chiral quark model [3, 4, 5, 6]. The latter implies the heavy-light $0^+, 1^+$ resonance states to form an anti-triplet representation of the SU(3) group. If one insists on a non-linear realization of the chiral SU(3) group only, excluding any further model assumptions, an a priori statement can not be made for the existence of chiral partners of any given state. Thus it is it is important to study the heavy-light meson and baryon resonances in great detail [7]

In this talk we review a particularly important aspect of open-charm resonances: the role of coupled-channel effects. It is long known that light scalar mesons can be described quantitatively in terms of coupled-channel dynamics without assuming the existence of light quark-antiquark states with scalar quantum numbers [8]. More recently it was pointed out by the authors [9] that similarly the light axial-vector meson spectrum of QCD is naturally explained in terms of chiral coupled-channel dynamics. The leading term of the chiral Lagrangian written down for the interaction of Goldstone bosons with light vector mesons generates the axial-vector spectrum dynamically. Moreover, chiral dynamics has the potential to predict the existence of exotic SU(3) multiplets. For instance in [10] weak attraction for s-wave scattering of Goldstone bosons off the decuplet ground states was found in the 27-plet channels. This may lead to the existence of penta-quark type resonance states in the $K\Delta$-channel. Since chiral coupled-channel dynamics is able to predict the s- and d-wave baryon resonance spectrum in terms of the leading order interaction term [10, 11] it is natural to expect coupled-channel dynamics to play also a crucial role in the physics of open-charm meson and baryon resonances. The possible importance of coupled-channel effects for the heavy-light meson resonances was

anticipated by Van Beveren and Rupp [12]. Systematic computations carried out by the authors [13, 14, 15] based on the chiral Lagrangian will be discussed here. For the possible role of diquark correlations in scalar open-charm meson resonances we refer to the talk of Terasaki [16, 17].

A coupled-channel effective field theory for the scattering of Goldstone bosons off heavy-light meson and baryon fields is presented [10, 11, 18, 19, 20, 21]. It relies on the chiral SU(3) Lagrangian with heavy-light $J^P = 0^-$, $J^P = 1^-$ and $J^P = \frac{1}{2}^+$ fields, that transform non-linearly under the chiral SU(3) group. The major result in the open-charm meson sector is the prediction that there exist chiral excitations in QCD with scalar and axial-vector quantum numbers forming anti-triplet and sextet representations of the SU(3) group. This differs from the results of the chiral quark model leading to anti-triplet states only [6]. It is pointed out that the heavy-quark symmetry, which suggests a degeneracy of scalar and axial-vector resonances at leading order, is naturally recovered in chiral coupled-channel theory. Further results are derived for open-charm baryon states. The s-wave scattering of Goldstone bosons off the anti-triplet and sextet ground states generates dynamically one anti-triplet, two sextet and one anti-quindecimet of states. This implies the existence of exotic penta-quark-like states with open-charm quantum numbers.

Particular results concern the 'heavy' SU(3) limit with $m_{\pi,\eta,K} \simeq 500$ MeV in which we predict bound states rather than resonance states typically. This observation should facilitate the verification of chiral coupled-channel dynamics via un-quenched QCD lattice simulations. In the 'light' SU(3) limit with $m_{\pi,\eta,K} \simeq 140$ MeV no clear resonance nor bound-state signals persist. The latter prediction poses a true challenge for un-quenched QCD lattice simulation, requiring small current quark masses as well as the analysis of extremely broad spectral distributions from Euclidian-space resonance propagators.
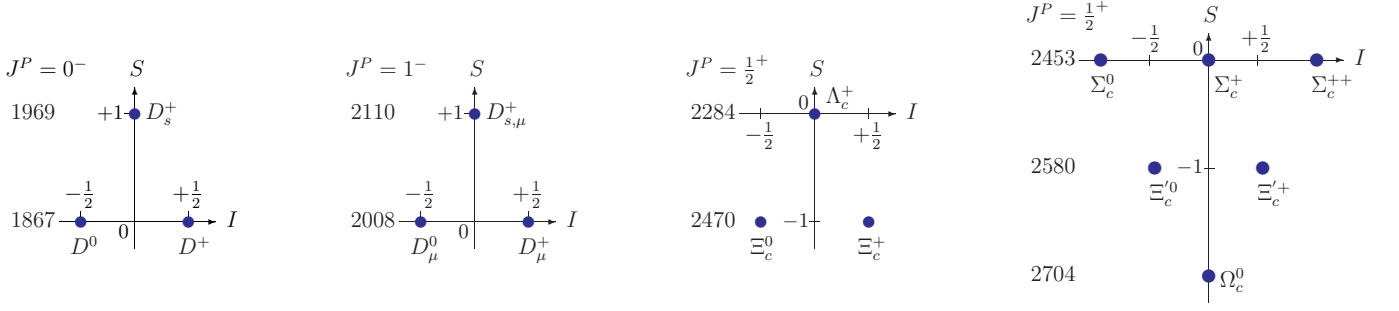
FIG. 1: Isospin (I) and strangeness (S) quantum numbers of open-charm mesons and baryons ground states with $J^P = 0^-, 1^-, \frac{1}{2}^+$. The isospin averaged masses in units of MeV of the various states are indicated.

II. CHIRAL COUPLED-CHANNEL DYNAMICS: THE χ-BS(3) APPROACH

The starting point of our study is the chiral SU(3) Lagrangian. We identify the leading-order term [22, 23, 24, 25] describing the interaction of Goldstone bosons with open-charm pseudo-scalar, vector mesons and spin-1/2 baryons,

$$
\begin{aligned}
\mathcal{L}(x) &= \frac{1}{8\,f^2}\,\mathrm{tr}\,\Big(\big(\big[P(x)\,(\partial^\nu P^\dagger(x)) - (\partial^\nu P(x))\,P^\dagger(x)\big] - \big[P^\mu(x)\,(\partial^\nu P^\dagger_\mu(x)) - (\partial^\nu P^\mu(x))\,P^\dagger_\mu(x)\big]\big)\,\big[\Phi(x), (\partial_\nu\,\Phi(x))\big]_-\Big) \\
&+ \frac{i}{16\,f^2}\,\mathrm{tr}\,\big(\bar{H}_{[\bar3]}(x)\,\gamma^\mu\,\big[H_{[\bar3]}(x),\,[\Phi(x),(\partial_\mu\,\Phi(x))]_-\big]_+\big) + \frac{i}{16\,f^2}\,\mathrm{tr}\,\big(\bar{H}_{[6]}(x)\,\gamma^\mu\,\big[H_{[6]}(x),\,[\Phi(x),(\partial_\mu\,\Phi(x))]_-\big]_+\big)\,. \quad (1)
\end{aligned}
$$

Here Φ is the Goldstone bosons field, P and P_μ are the massive pseudo-scalar and vector-meson fields and and $H_{[\bar3]}$ and $H_{[6]}$ the massive baryon fields. The $J^P = 0^-$ and $J^P = 1^-$ fields, P and P_μ, form anti-triplet representations of the SU(3) flavour group. The $J^P = \frac{1}{2}^+$ baryon fields $H_{[\bar3]}$ and $H_{[6]}$ transform as anti-triplet and sextet. The corresponding isospin (I) and strangeness (S) quantum numbers are recalled in Fig. 1. Finally the parameter f in (1) is known from the weak decay process

of the pions. We use $f = 90$ MeV.

The scattering problem decouples into seven orthogonal channels specified by isospin and strangeness quantum numbers. Heavy-light meson and baryon resonances with quantum numbers $J^P = 0^+$, $J^P = 1^+$ and $J^P = \frac{1}{2}^-$ manifest themselves as poles in the s-wave scattering amplitudes, $M_{J^P}^{(I,S)}(\sqrt{s})$, which in the $\chi-$BS(3) approach [9, 21] take the simple form

$$
M_{J^P}^{(I,S)}(\sqrt{s}) = \Big[1 - V_{J^P}^{(I,S)}(\sqrt{s})\,J_{J^P}^{(I,S)}(\sqrt{s})\Big]^{-1} V_{J^P}^{(I,S)}(\sqrt{s})\,, \qquad V_{\frac{1}{2}^+}^{(I,S)}(\sqrt{s}) = \frac{C_{\frac{1}{2}^+}^{(I,S)}}{4\,f^2}\Big(2\,\sqrt{s} - M - \bar{M}\Big)\,,
$$

$$
V_{0^+}^{(I,S)}(\sqrt{s}) = V_{1^+}^{(I,S)}(\sqrt{s}) = \frac{C_{0^+}^{(I,S)}}{8\,f^2}\Big(3\,s - M^2 - \bar{M}^2 - m^2 - \bar{m}^2 - \frac{M^2 - m^2}{s}\,(\bar{M}^2 - \bar{m}^2)\Big)\,, \qquad (2)
$$

where (m, M) and $(\bar{m}, \bar{M})$ are the masses of initial and final states. We use capital M for the masses of heavy-light fields and small m for the masses of the Goldstone bosons. The effective interaction kernel $V^{(I,S)}(\sqrt{s})$ in (2) is determined by the leading-order chiral SU(3) Lagrangian (1).

The enormous predictive power of the chiral Lagrangian lies in its specification of the coupled-channel

matrices $C_{J^P}^{(I,S)}$. For the scattering of Goldstone bosons off any SU(3) anti-triplet field the latter may be classified by invariants according to the decomposition

$$
\bar3 \otimes 8 = \bar3 \oplus 6 \oplus \overline{15}\,. \qquad (3)
$$

The leading order chiral Lagrangian predicts attraction in the anti-triplet and sextet channels but repulsion for the anti-quindecimet. This holds irrespective whether

the target is a meson or baryon. What matters is that the target transforms as an anti-triplet under the SU(3) flavour group. A further prediction of chiral SU(3) symmetry is that the attraction in the anti-triplet is thrice as strong as the attraction in the sextet sector. A similar analysis reveals that the scattering of Goldstone bosons off sextet targets

$$6 \otimes 8 = \bar{3} \oplus 6 \oplus \overline{15} \oplus 24 \,, \qquad (4)$$

is attractive in the anti-triplet, sextet but also in the anti-quindecimet channels. The interaction is repulsive in the 24-plet channels. Chiral SU(3) symmetry predicts a hierarchy of strength with strongest attraction in the triplet channels, which is five times as strong as the attraction

in the anti-quindecimet channels. In the sextet channel the attraction is reduced by a factor 3/5 only. It is also instructive to compare the amount of attraction in the anti-triplet channels as they result from the reduction of $\bar{3} \otimes 8$ versus $6 \otimes 8$. Chiral symmetry predicts stronger binding in the latter case. The amount of attraction is larger by a factor 5/3 as compared to the former case [31].

It is left to specify the loop functions $J_{J^P}^{(I,S)}(\sqrt{s})$ in (2). The latter are diagonal in the coupled-channel space, depend however on whether to scatter Goldstone bosons off mesons or baryons. For details we refer to [9, 13, 15, 21]. The final expressions are quite transparent,

$$J_{0^+}(\sqrt{s}) = I(\sqrt{s}) - I(\mu_{0^+}) \,, \qquad J_{1^+}(\sqrt{s}) = \left(1 + \frac{p^2}{3\,M^2}\right)\left[I(\sqrt{s}) - I(\mu_{1^+})\right] \,, \qquad (5)$$

$$J_{\frac{1}{2}^+}(\sqrt{s}) = \left(M + (M^2 + p^2)^{1/2}\right)\left[I(\sqrt{s}) - I(\mu_{\frac{1}{2}^+})\right] \,,$$

$$I(\sqrt{s}) = \frac{1}{16\,\pi^2}\left(\frac{p}{\sqrt{s}}\left(\ln\left(1 - \frac{s - 2\,p\,\sqrt{s}}{m^2 + M^2}\right) - \ln\left(1 - \frac{s + 2\,p\sqrt{s}}{m^2 + M^2}\right)\right) + \left(\frac{1}{2}\frac{m^2 + M^2}{m^2 - M^2} - \frac{m^2 - M^2}{2\,s}\right)\ln\left(\frac{m^2}{M^2}\right)\right) \,,$$

where $\sqrt{s} = \sqrt{M^2 + p^2} + \sqrt{m^2 + p^2}$. Note that the loop functions in (6) specifying the scattering of Goldstone bosons off 0^- and 1^- targets differ by a term suppressed with $1/M^2$ only. A crucial ingredient of the χ−BS(3) scheme is its approximate crossing symmetry guaranteed by a proper choice of the subtraction scales,

$$\begin{aligned}
\mu_{0^+}^{(I,0)} &= M_{D(1867)} \,, & \mu_{0^+}^{(I,\pm 1)} &= M_{D_s(1969)} \,, \\
\mu_{0^+}^{(I,2)} &= M_{D(1867)} \,, & \mu_{1^+}^{(I,0)} &= M_{D(2008)} \,, \\
\mu_{1^+}^{(I,\pm 1)} &= M_{D_s(2110)} \,, & \mu_{1^+}^{(I,2)} &= M_{D(2008)} \\
\mu_{[3]}^{(I,0)} &= M_{\Lambda_c(2284)} \,, & \mu_{[3]}^{(I,\pm 1)} &= M_{\Xi_c(2470)} \,, \quad (6) \\
\mu_{[3]}^{(I,-2)} &= M_{\Lambda_c(2284)} \,, & \mu_{[6]}^{(I,0)} &= M_{\Sigma_c(2453)} \,, \\
\mu_{[6]}^{(I,\pm 1)} &= M_{\Xi_c'(2580)} \,, & \mu_{[6]}^{(I,-2)} &= M_{\Omega_c(2704)} \,, \\
\mu_{[6]}^{(I,-3)} &= M_{\Xi_c'(2580)} \,.
\end{aligned}$$

The renormalization condition (6) reflects the basic assumption our effective field theory is built on, namely, that at subthreshold energies the scattering amplitudes may be evaluated in standard chiral perturbation theory. Once the available energy is sufficiently high to permit elastic two-body scattering, a further typical dimensionless parameter of order one arises, that is uniquely linked to the presence of a two-particle unitarity cut. Thus it is sufficient to sum those contributions keeping the perturbative expansion of all terms that do not develop a two-particle unitarity cut. In order to recover the perturbative nature of the subthreshold scattering amplitude the subtraction scale $M - m < \mu < M + m$ must be chosen in between the s- and u-channel elastic unitarity

branch points [21]. It was suggested that s-channel and u-channel unitarized amplitudes should be glued together at subthreshold kinematics [21]. A smooth result is guaranteed if the full amplitudes match the interaction kernel V close to the subtraction scale μ as implemented by (6). In this case the crossing symmetry of the interaction kernel, which follows directly from its perturbative evaluation, is carried over to an approximate crossing symmetry of the full scattering amplitude. This construction reflects our basic assumption that diagrams showing an
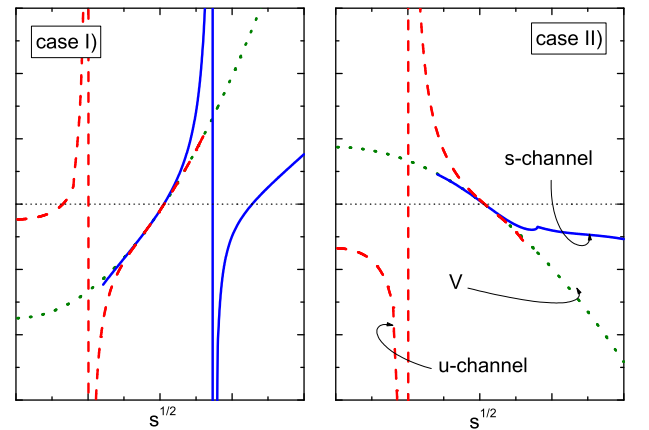

FIG. 2: Typical cases of forward scattering amplitudes. The solid (dashed) line shows the s-channel (u-channel) unitarized scattering amplitude. The dotted lines represent the amplitude evaluated at tree-level.

s-channel or u-channel unitarity cut need to be summed to all orders typically at energies where the diagrams develop their imaginary part. In Fig. 2 we demonstrate the quality of the proposed matching procedure as applied for typical forward scattering amplitudes. The figure clearly illustrates the smooth matching of s-channel and u-channel iterated amplitudes at subthreshold energies.

III. DISCUSSION OF RESULTS

We first discuss the resonance spectrum as it arises in the 'heavy' SU(3) limit [9, 10] with $m_{\pi, K, \eta} = 500$ MeV. Using degenerate masses of 1800 MeV for the pseudo-scalar and vector D-mesons an almost degenerate spectrum of scalar and axial-vector D-mesons is obtained. The mass splitting are below 1 MeV reflecting the presence of the heavy-quark symmetry. In this case an anti-triplet of mass 2204 MeV with poles in the $(I, S) = (0, +1), (1/2, 0)$ amplitudes is generated dynamically. In the sextet channels the attraction is not quite strong enough to form a bound or clear resonance state signal. However, if the attraction is increased slightly by using $f = 80$ MeV rather than the canonical value 90 MeV, poles at mass 2298 MeV arise in the $(1, +1), (1/2, 0), (0, -1)$ amplitudes. Similar results are produced for the s-wave open-charm baryon resonances. Using a somewhat arbitrary common mass 2400 MeV for the anti-triplet ground states an anti-triplet of mass 2778 MeV with poles in the $(0, 0), (\frac{1}{2}, -1)$ amplitudes and a sextet of mass 2900 MeV with poles in the $(1, 0), (\frac{1}{2}, -1), (0, -2)$ amplitudes is generated dynamically. The result is quite stable against small variations of the optimal subtraction scales of (6). Lowering the latter by 200 MeV reduces the anti-triplet and sextet masses by 40 MeV and 5 MeV only. Assuming a common mass for the sextet ground states of 2500 MeV we obtain an anti-triplet of mass 2807 MeV with poles in the $(0, 0), (\frac{1}{2}, -1)$ amplitudes, a sextet of mass 2875 MeV with poles in the $(1, 0), (\frac{1}{2}, -1), (0, -2)$ amplitudes and an anti-quindecimet of mass 3000 MeV with poles in the $((\frac{1}{2}, 1), (0, 0), (1, 0), (\frac{1}{2}, -1), (\frac{3}{2}, -1), (1, -2))$ amplitudes.

Contrasted results follow in the 'light' SU(3) limit [9, 10] with $m_{\pi, K, \eta} \sim 140$ MeV. No clear signal of a resonance or bound state in any of the channels is found. The striking dependence of the resonance spectrum on the current quark masses of QCD should be tested in un-quenched lattice simulations.

Realistic results with clear bound-state or resonance signals in many channels are implied by using physical masses of the ground state open-charm meson and baryons. We first discuss the open-charm meson spectrum. In the $(0, 1)$-sector we predict a scalar bound state of mass 2303 MeV and an axial-vector state of mass 2440 MeV at leading chiral order. According to [5, 6] these states should be identified with the narrow resonances of mass 2317 MeV and 2463 MeV recently observed by

the BABAR collaboration [1, 2]. Since we do not consider isospin violating processes like $\eta \to \pi_0$ the latter states are true bound state in our present scheme. Given the fact that our computation is parameter-free this is a remarkable result.

Further clear scalar and axial-vector resonance signals are seen in the $(\frac{1}{2}, 0)$-sector at leading order. Chiral dynamics predicts a narrow scalar state of mass 2413 MeV just below the $\eta D(1867)$-threshold and a broad scalar state of mass 2138 MeV. Modulo some mixing effects the heavier of the two is part of the sextet the lighter a member of the anti-triplet. Heavy-quark symmetry implies an analogous pattern for the axial-vector resonances. A narrow structure at 2552 MeV arises together with a broad state at around 2325 MeV. The two broad states, which were first predicted by the chiral quark model [6, 12], were confirmed by the recent data of the BELLE collaboration [26]. In Fig. 3 we display the data together with the result of [14] which performed a chiral coupled-channel computation at subleading order. In the right-hand panel three resonances contribute to the measured $\pi D(2008)$-spectrum. The effect of a narrow 2^+ state of mass 2460 MeV indicated by the histogram is small and not considered in the computation [14]. The remaining two states are identified with a broad anti-triplet and a narrow sextet state. This implies that the well established $D(2420)$-resonance [27] should be a member of a sextet rather than a member of an anti-triplet as commonly assumed. Of course this questions also the in-
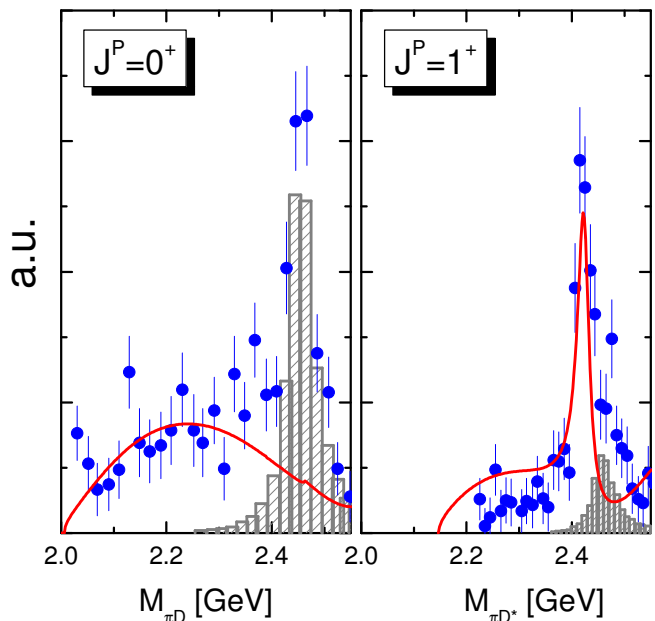

FIG. 3: Mass spectra of the $(\frac{1}{2}, 0)$-resonances as seen in the $\pi D(1867)$-channel (l.h. panel $J^P = 0^+$) and $\pi D(2008)$-channel (r.h. panel $J^P = 1^+$). The solid lines show the theoretical mass distributions. The data are taken from [26] as obtained from the $B \to \pi D(1867)$ and $B \to \pi D(2008)$ decays. The histograms indicate the contribution from the $J = 2$ resonances $D(2460)$ as given in [26].
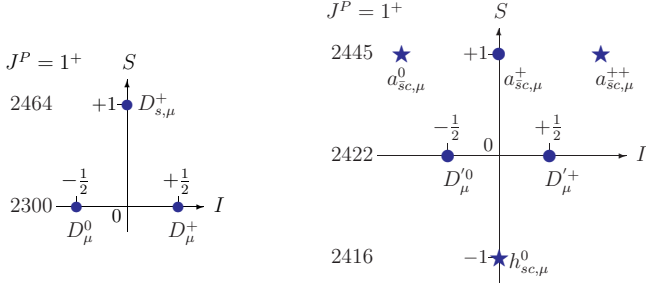
FIG. 4: Isospin (I) and strangeness (S) quantum numbers of open-charm mesons with $J^P = 1^+$. Masses in units of MeV of the chiral excitations are indicated. Resonances with exotic quantum numbers are represented by stars.

terpretation of the 2^+ state of mass 2460 MeV to be the heavy-quark partner of the $D(2420)$ suggesting at least one additional so far unobserved state, possibly with quantum numbers 3^+. To finally settle this issue requires certainly further work. It should be stressed that in the computation [14] free parameters which enter at subleading order were adjusted to recover the $D(2420)$ resonance at its physical mass. At leading order the mass of the latter is overestimated by about 130 MeV even though a small width of about 20 MeV is obtained. Thus the chiral correction terms increase the amount of attraction in the sextet channel leading to the level pattern shown in Fig. 4. In particular a $KD(2008)$-molecule state of mass 2416 MeV, which carries exotic quantum numbers, is predicted. We suggest to call the latter h_{sc} in analogy to the axial-vector isospin-singlet state $h(1170)$. Similarly, the name $a_{\bar{s}c}$ for the isospin-triplet members of the sextet is proposed. The various reduced scattering amplitudes are shown in channels where the generated states have a finite width are shown in Fig. 5. The figure clearly illustrates the typical phenomenon observed in chiral coupled-channel theory: the $(I,S) = (1/2,0)$-state couples most strongly to the kinematically closed channels $\eta D(2008)$ and $\bar{K}D_s(2110)$.

A spectrum similar to the one for the axial-vector mesons arises for the open-charm scalar mesons. At leading order two isospin-doublet resonances with zero strangeness are generated by coupled-channel effects. The triplet state couples strongly to the $\pi D(1868)$-channel leading to its large width. The sextet state couples only weakly to the latter channel implying a narrow state. In Fig. 3 an extreme scenario is shown in which the sextet state decouples from the $\pi D(1868)$-channel all together. This can be achieved upon considering chiral correction terms [14]. The sextet state with zero strangeness sits at 2389 MeV in this case. The $KD(1867)$ molecule state comes at 2352 MeV. The signal of the sextet is weakest in the $(I,S) = (1,1)$ sector where the scattering amplitude shows a strong cusp effect close to the $KD(1867)$-threshold only. We emphasize that it is difficult to make a precise prediction for the width of the iso-scalar sextet state. If we only slightly change the set of parameters the sextet state shows up as a narrow peak
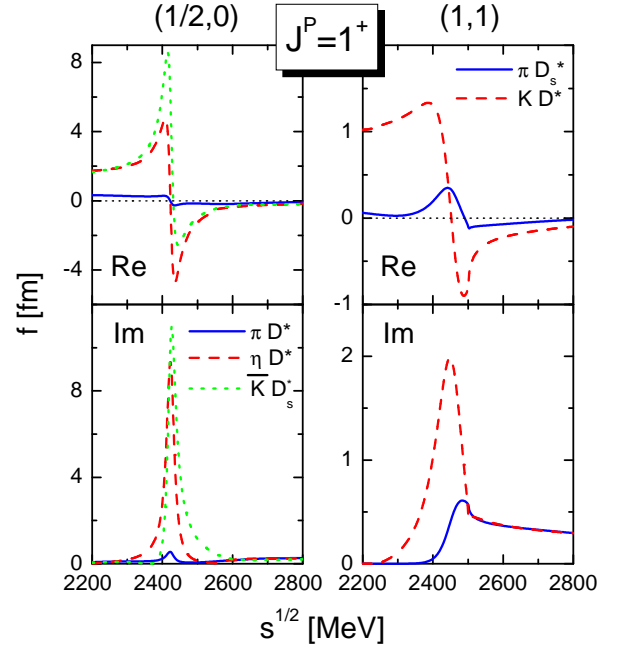

FIG. 5: Open-charm resonances with $J^P = 1^+$ and $(I,S) = (\frac{1}{2},0), (0,1)$ as seen in the scattering of Goldstone bosons of $D(2008)$- and $D_s(2110)$-mesons. Shown are real and imaginary parts of reduced scattering amplitudes.

in the $\pi D(1867)$-spectrum of Fig. 3. Depending on the precise values of the parameters this state may be detected most efficiently via its coupling to the $\eta D(1867)$ channel utilizing the $\eta - \pi_0$-mixing effect. In any case it is of utmost importance to further improve the quality of the empirical spectrum. The existence of scalar open-charm mesons with exotic quantum numbers was suggested first by Terasaki based on a phenomenological constituent diquark approach [16].

We turn to the chiral excitations of open-charm baryons. Unfortunately, at present there is very little known empirically about the open-charm baryon resonance spectrum. Only three states $\Lambda_{c1}(2593), \Lambda_c(2880)$ and $\Xi_c(2790)$ were discovered by the CLEO collaboration so far [27]. The level scheme, as it is predicted by chiral dynamics at leading order using physical values
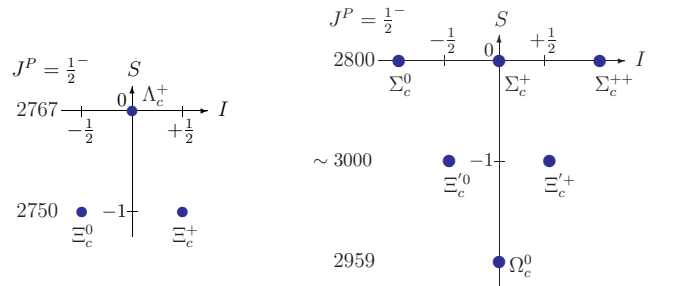

FIG. 6: Isospin (I) and strangeness (S) quantum numbers of open-charm baryons with $J^P = \frac{1}{2}^-$ generated dynamically by scattering Goldstone bosons off the triplet ground states. Masses in units of MeV of the chiral excitations are indicated.

for the ground state masses, is drawn in Fig. 6 for the s-wave resonances generated by the scattering of Goldstone bosons off the anti-triplet ground states. We obtain a bound state of mass 2767 MeV and quantum numbers $(I, S) = (0, 0)$. This state should be identified with the $\Lambda_c(2880)$ recently detected by the CLEO collaboration [28] via its decay into the $\pi\Sigma_c(2453) \to \Lambda_c\pi\pi$ channel. The narrow width of the observed state of smaller than 8 MeV [28] appears consistent with a suppressed coupling of that state to the $\pi\Sigma_c(2453)$ channel as predicted by chiral symmetry. The various scattering amplitudes in the $(0, 0)$ and $(1/2, -1)$ sectors are collected in Fig. 7. Similarly the scattering amplitudes [32] probing the sextet states of Fig. 6 are collected in Fig. 8. The figures clearly indicate that it is much too simple showing a schematic level diagram only. The amplitudes exhibit a complicated multi-channel structure. For instance whereas the scattering amplitudes show a clear signal for the SU(3) triplet state $\Xi_c(2750)$ a corresponding state belonging to the sextet is quite broad and does not manifest itself very clearly in the amplitudes of Fig. 7. The existence of the sextet may be best confirmed searching for an isospin-singlet $K\,\Xi_c(2470)$-molecule.

We conclude our discussion with the spectrum generated dynamically by scattering Goldstone bosons off the sextet ground states. The resulting level scheme is drawn schematically in Fig. 9. Since the interaction is attractive in the anti-triplet, sextet and anti-quindecimet channels a quite rich spectrum arises. The figure suggests a particular labeling of the exotic multiplets. In particular we denote an isospin-triplet state with strangeness minus-two by Σ_{ssc}. Similarly we introduce the isospin-quartet, Δ_{sc}, which carries strangeness minus-one. Again the level scheme has to be taken with great care since not in all multiplet channels clear resonance structures are expected. It is an important part of the predictive power of chiral coupled-channel dynamics to spell out in which channels one expects clear signals. Full details of the resonance properties are provided by Figs. 7-10 in terms of reduced scattering amplitudes. For the channels in which the chiral excitations of the anti-triplet and sextet manifest themselves with states of identical isospin and strangeness quantum numbers the corresponding amplitudes are grouped together (see Figs. 7-8). This is instructive since at subleading order the amplitudes start to couple.
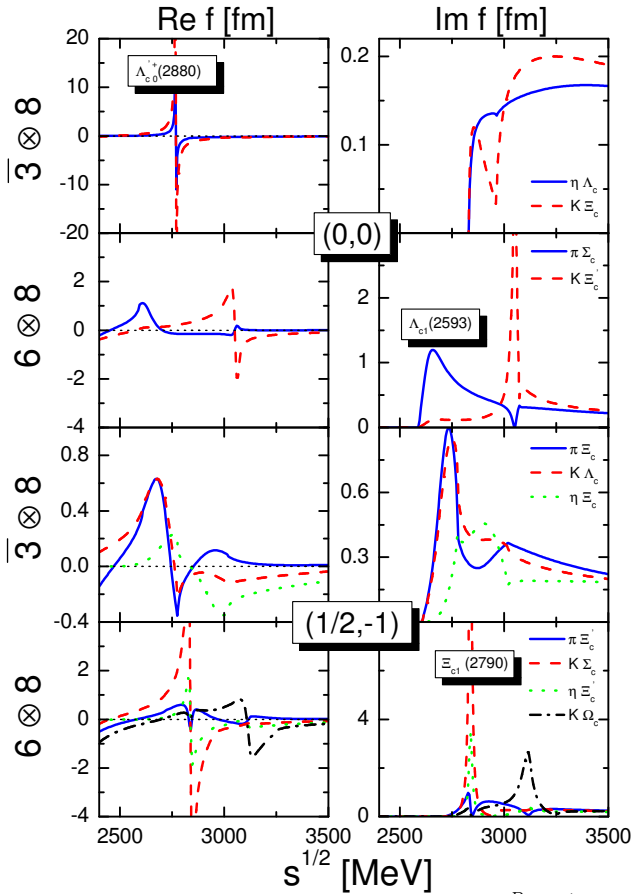

FIG. 7: Open-charm baryon resonances with $J^P = \frac{1}{2}^-$ and $(I, S) = (0, 0)$ and $(1/2, -1)$ as seen in the scattering of Goldstone bosons off anti-triplet $(\Lambda_c(2284), \Xi_c(2470))$ and sextet $(\Sigma_c(2453), \Xi'_c(2580), \Omega_c(2704))$ baryons. Shown are real and imaginary parts of reduced scattering amplitudes.
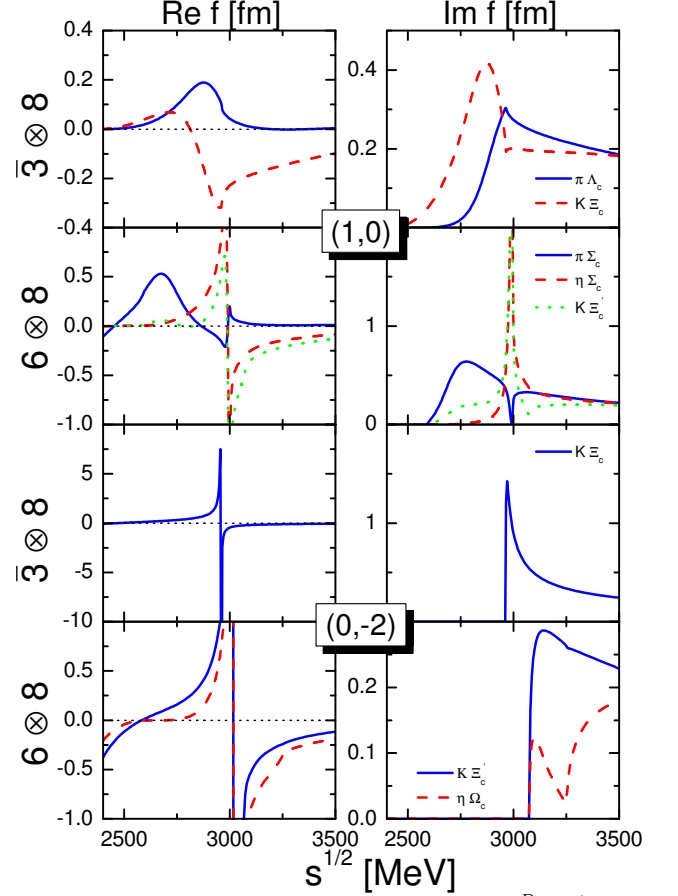

FIG. 8: Open-charm baryon resonances with $J^P = \frac{1}{2}^-$ and $(I, S) = (1, 0)$ and $(0, -2)$ as seen in the scattering of Goldstone bosons off anti-triplet $(\Lambda_c(2284), \Xi_c(2470))$ and sextet $(\Sigma_c(2453), \Xi'_c(2580), \Omega_c(2704))$ baryons. Shown are real and imaginary parts of reduced scattering amplitudes.
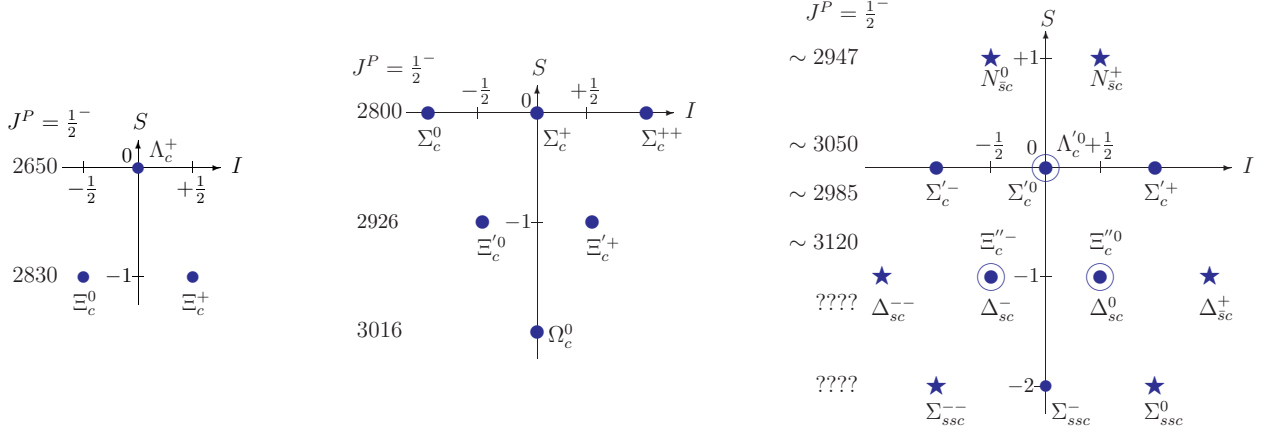
FIG. 9: Isospin (I) and strangeness (S) quantum numbers of open-charm baryons with $J^P = \frac{1}{2}^-$ generated dynamically by scattering Goldstone bosons off the sextet ground states. Masses in units of MeV of the chiral excitations are indicated. Resonances with exotic quantum numbers are represented by stars.

Consider first the anti-triplet states that are generated dynamically as chiral excitations of the sextet ground states. Since the attraction is largest in this channel one expects strongly bound systems with $(I, S) = (0, 0)$ and $(1/2, -1)$. Indeed the second row of Fig. 7 shows a resonance at about 2650 MeV that couples strongly to the $\pi\Sigma_c(2453)$ channel. The properties of this state are close to the ones of the $\Lambda_c(2593)$ resonance [29]. Here we obtain a decay width which is significantly larger than the empirical width of about 4 MeV [29]. Chiral correction terms that couple the states seen in the 1st and 2nd rows of Fig. 7 are expected to decrease this width. Level-

level repulsion of the two observed states should lower the mass of the lighter state but push up the mass of the heavier state. In the 4th row of Fig. 7 a narrow state at 2830 MeV is shown. This state couples strongly to the $K\Sigma_c$ and $\eta\Xi_c$ channels and due to its small width should be identified with the $\Xi_c(2790)$ resonance [30].

The sextet manifests itself with a bound-state of mass 3015 MeV with $(I, S) = (0, -2)$ as illustrated in the 4th row of Fig. 8. This channel is a unique probe for the sextet since neither the ant-triplet nor the anti-quindecimet contributes here. A further broad sextet state of mass ~ 2800 MeV with $(1, 0)$ is seen in the 2nd row of Fig. 8. It couples strongly to the $\pi\Sigma_c(2453)$-channel. Finally the second broad resonance of mass 2926 MeV in the 4th row of Fig. 7 that couples strongly to the $\pi\Xi_c(2580)$-channel should also be a member of the sextet.

The most exciting consequence of chiral coupled-channel dynamics is its prediction of anti-quindecimet states. As indicated in Fig. 9 we do not expect clear signals in all channels. Most prominent are the states with $(I, S) = (0, 0), (1, 0), (1/2, -1)$ around 3 GeV as shown in the 2nd rows of Figs. 7-8 and 4th row of Fig. 7. Unfortunately these states do not carry exotic quantum numbers. The channels in which the anti-quindecimet manifests itself with exotic quantum numbers are displayed in Fig. 10. One may hope that chiral corrections terms conspire to further increase the amount of attraction leading to signals that are easier to observe.
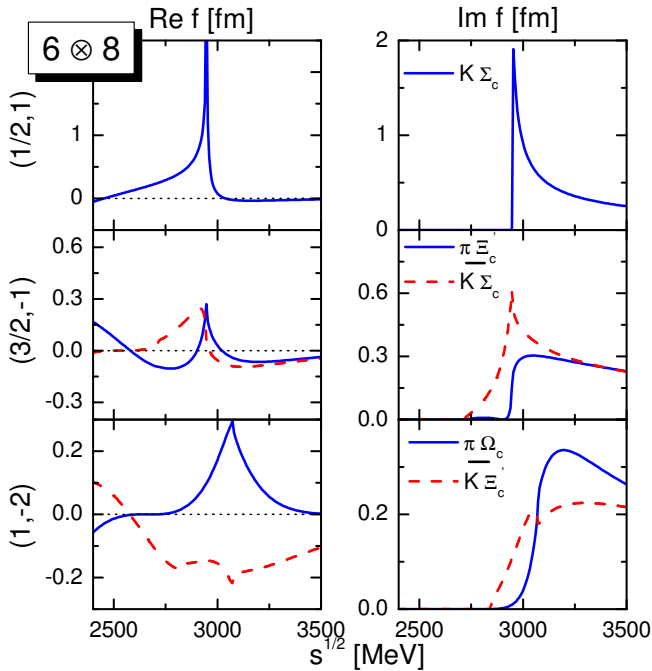

FIG. 10: Open-charm baryon resonances with $(I, S) = (\frac{1}{2}, 1), (\frac{3}{2}, -1), (1, -2)$ and $J^P = \frac{1}{2}^+$ as seen in the scattering of Goldstone bosons off sextet $(\Sigma_c(2453), \Xi_c(2580), \Omega_c(2704))$ baryons. Shown are real and imaginary parts of reduced scattering amplitudes.

IV. SUMMARY

We reviewed the spectrum of chiral excitations in the open-charm meson and baryon sector. Exciting new predictions for the existence of exotic multiplets were presented in detail. These results suggest a dedicated search for such states but also require extensive chiral coupled-channel computations to further substantiate the theoretical predictions.

[1] BABAR Collaboration, B. Aubert et al., Phys. Rev. Lett. **90** (2003) 242001.

[2] CLEO Collaboration, D. Besson et al., Phys. Rev. **D 68** (2003) 032002; hep-ex/0305017.

[3] W.A. Bardeen and C.T. Hill, Phys. Rev. **D 49** (1994) 409.

[4] M.A. Nowak, M. Rho and I. Zahed, Phys. Rev. **D 48** (1993) 4370.

[5] W.A. Bardeen, E.J. Eichten and Ch.T. Hill, Phys. Rev. **D 68** (2003) 054024.

[6] M.A. Nowak, M. Rho and I. Zahed, hep-ph/0307102.

[7] R.N. Cahn and J.D. Jackson, Phys. Rev. **D 68** (2003) 037502; T. Barnes, F.E. Close and H.J. Lipkin, Phys. Rev. **D 68** (2003) 054006; H.-Y. Cheng and W.-S. Hou, Phys. Lett. **B 566** (2003) 193; A.P. Szczepaniak, Phys. Lett. **B 567** (2003) 23; S. Godfrey, Phys. Lett. **B 568** (2003) 254; P. Colangelo and F. De Fazio, Phys. Lett. **B 570** (2003) 180; G. S. Bali, Phys. Rev. **D 68** (2003) 071501; S. Nussinov, hep-ph/0306187; Y.-B. Dai, C.-S. Huang, C. Liu and S.-L. Zhu, Phys. Rev. **D 68** (2003) 114011; A. Dougall, R.D. Kenway, C.M. Maynard and C. McNelie, Phys. Lett. **B 569** (2003) 41; Th.E. Browder, S. Pakvasa and A.A. Petrov, Phys. Lett. **B 578** (2004) 365; A. Deandrea, G. Nardulli, A.D. Polosa, Phys. Rev. **D 68** (2003) 097501; Ch.-H. Chen and H.N. Li, Phys. Rev. **D 69** (2004) 054002; M. Sadzikowski, Phys. Lett. **B 579** (2004) 39; A. Datta and P.J. O'donnell, Phys. Lett. **B 572** (2003) 164; M. Suzuki, hep-ph/0307118; P. Bicudo, hep-ph/0401106.

[8] E. Van Beveren et al., Z. Phys. **C 30** (1986) 615; J.D. Weinstein and N. Isgur, Phys. Rev. **D 41** (1990) 2236; G. Janssen, B. C. Pearce, K. Holinde, J. Speth, Phys. Rev. **D 52** (1995) 2690; J.A. Oller, E. Oset and J.R. Pelaez, Phys. Rev. **D 59** (1999) 074001; Erratum-ibid. **D 60** (1999) 099906; J. Nieves, M.P. Valderrama and E.R. Arriola, Phys. Rev. **D 65** (2002) 036002; A.G. Nicola and J.R. Pelaez, Phys. Rev. **D 65** (2002) 054009;

[9] M.F.M. Lutz and E. E. Kolomeitsev, Nucl. Phys. **A 730** (2004) 392.

[10] E.E. Kolomeitsev and M.F.M. Lutz, Phys. Lett. **B 585** (2004) 243.

[11] C. García-Recio, M.F.M. Lutz and J. Nieves, Phys. Lett. **B 582** (2004) 49.

[12] E. van Beveren and G. Rupp, Phys. Rev. Lett. **91** (2003) 012003; Eur. Phys. J. **C 32** (2004) 493.

[13] E.E. Kolomeitsev und M.F.M. Lutz, Phys. Lett. **B 582** (2004) 39.

[14] J. Hofmann and M.F.M. Lutz, Nucl. Phys. **A 733** (2004) 142.

[15] M.F.M. Lutz and E.E. Kolomeitsev, Nucl. Phys. **A 730** (2004) 110.

[16] K. Terasaki, Phys. Rev. **D 68** (2003) 011501.

[17] K. Terasaki, help-ph/0405146.

[18] M.F.M. Lutz and E. E. Kolomeitsev, Proc. of Int. Workshop XXVIII on Gross Properties of Nuclei and Nuclear Excitations, Hirschegg, Austria, January 16-22, 2000.

[19] M.F.M. Lutz und E.E. Kolomeitsev, Found. Phys. **31** (2001) 1671.

[20] M.F.M. Lutz, GSI-Habil-2002-1.

[21] M.F.M. Lutz and E. E. Kolomeitsev, Nucl. Phys. **A 700** (2002) 193.

[22] S. Weinberg, Phys. Rev. Lett. **17** (1966) 616; Y. Tomozawa, Nuov. Cim. **A 46** (1966) 707.

[23] M.B. Wise, Phys. Rev. **D 45** (1992) 2188.

[24] T.-M. Yan et al., Phys. Rev. **46** (1992) 1148.

[25] G. Burdman and J. Donoghue, Phys. Lett. **B 280** (1992) 287.

[26] BELLE Collaboration, hep-ex/0307021.

[27] K. Hagiwara et al., Phys. Rev. **D 66** (2002) 010001.

[28] CLEO Collaboration, M. Artuso et. al., Phys. Rev. Lett. **86** (2001) 4479.

[29] CLEO Collaboration, K.W. Edwards et al., Phys. Rev. Lett. **74** (1995) 3331.

[30] CLEO Collaboration, S.E. Csorna et. al., Phys. Rev. Lett. **86** (2001) 4243.

[31] We correct two minor misprints in Tab. 3 of [15]. The 11 components of the second and fifth row should read -2 rather than -1 and 0.

[32] The second row corrects a slightly incorrect figure shown in [15].